

Freed Leptogenesis

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April 24, 2012

Economical extensions of the Standard Model (SM), in which famous Davidson–Ibarra bound on the CP asymmetry relevant for leptogenesis may be significantly relaxed by the loop effects, comparing to predictions of the SM extended only by heavy right-handed neutrinos with hierarchical masses, are discussed. This leads to decreasing of the lower bound on the heavy neutrino masses and increasing of the upper bound on the light neutrino masses, which is testable. In addition, the considered theory may help to solve the dark matter problem.

1 Introduction

The observable small nonzero neutrino masses and baryon asymmetry of the Universe (BAU) [1] provide strong evidences of physics beyond the Standard Model (SM). The see-saw mechanism [2, 3, 4, 5, 6, 7] gives economical explanation of the lightness of neutrinos by adding the heavy Majorana neutrinos to the SM particle content, which generate the small neutrino masses by the tree level perturbative interaction with the SM Higgs vacuum expectation value (VEV). In addition, the BAU may be explained by generating the lepton asymmetry in the out-of-equilibrium decays of these heavy neutrinos and converting it to a baryon asymmetry by sphaleron transitions [8] in the usual baryogenesis [9] via leptogenesis (LG) scenario [10]. However the successful LG in this simple SM extension requires (in the case of hierarchical heavy neutrinos) a strong upper bound on the relevant CP asymmetry, which was introduced in [11, 12, 13, 14], and is called Davidson–Ibarra (DI) bound. This results also in the lower bound on the right-handed neutrino masses of $\sim 10^9$ GeV [14] and the upper bound on the left-handed neutrino masses of ~ 0.1 eV [15, 16]. By generalizing the SM to the Minimal Supersymmetric Model the bound on the CP asymmetry is increasing only by the factor of two, which leads to the famous gravitino problem [17, 18, 19]. However in the case of quasi-degenerate heavy neutrinos a resonant enhancement of the CP asymmetry may happen [20, 21].

Another possible solution for the problem of small observable neutrino masses is its radiative generation [22, 23, 24, 25, 26, 27, 28, 29]. In this paper we consider generation

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of the neutrino masses at both tree and loop levels. We show that the theories with analytical relation between the couplings relevant for tree and loop contributions to the neutrino masses may significantly relax the DI bound in the case when these tree and loop terms approximately cancel each other. As a result, strongly hierarchical heavy neutrino masses in this theory may be tested at the Large Hadron Collider (LHC) and next particle facilities [30, 31, 32, 33, 34, 35, 36]. The discussed analytical relation may come from the structure of grand unified theories (GUT) [37], in which the particles involved in the tree and loop contributions to the neutrino masses belong to the same multiplets. In particular, Renormalizable Adjoint $SU(5)$ model [38] is one of the minimal realistic GUTs, in which a linear relation of this type is realized [39].

In the next section we investigate generation of the neutrino masses in the economical SM extensions with the loop contribution to the neutrino masses analogous to [26] and [27, 28, 29]. We analyze LG in section 3, and conclude in section 4.

2 Generation of neutrino masses

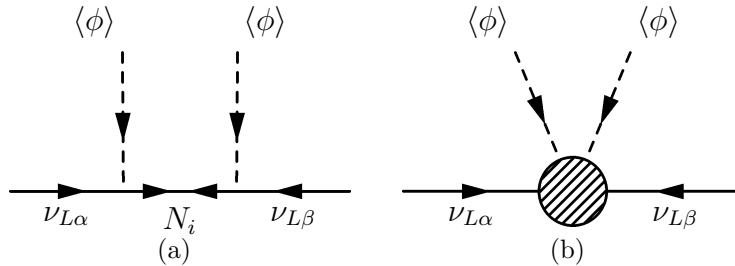


Figure 1: Considered contributions to the neutrino masses. The arrows in fermionic lines show flow of lepton number.

Consider theory with the neutrino masses generated by heavy Majorana fermions N_i ,¹ as shown in Fig. 1 *a*, and by other new heavy particles, which is shown effectively in Fig. 1 *b* after integration out these particles. It is well known that besides generating the neutrino masses the heavy fermions N_i can be at the same time responsible for the LG. In the case when the new heavy particles, involved in the contribution in Fig. 1 *b*, are decoupled from LG this contribution may relax the connection between the neutrino masses and LG, namely the DI bound. Such *LG* we call *Freed*.

In this paper we discuss a particular class of theories with the dominant 1-loop contribution to the effective vertex in Fig. 1 *b*, shown in Fig. 2, where N is new $SU(2)_L$ singlet Majorana fermion and η is new $SU(2)_L$ doublet scalar. The two possible classes of models, which generate this contribution, were introduced by Ma [26] and Perez–Wise [27, 28, 29]. Consider extensions of this models by several singlet or singlet

¹The correspondent mechanism of generation of the neutrino masses is called type I or type III see-saw, depending on whether N_i is singlet or triplet fermion, respectively.

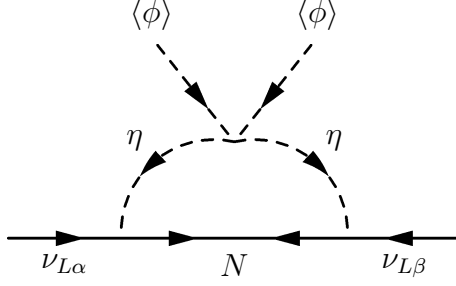


Figure 2: Possible 1-loop contribution to neutrino masses. The arrows in fermionic lines show flow of lepton number.

and triplet Majorana fermions N_i . In the minimal case we need only two N_i for non-degenerate neutrino masses and successful LG. The new particles in these extended Ma (EMM) and Perez–Wise (EPWM) models with their properties under the SM groups $G_{\text{SM}} = SU(3)_c \times SU(2)_L \times U(1)_Y$ and the discrete symmetry Z_2 (in the case of EMM) are listed in Tables 1 and 2, respectively. All the SM particles in EMM have positive Z_2 parity. We take this definitions of extended models for two reasons: simplicity in the case of EMM, and reproduction of particles responsible for the neutrino masses and LG of Adjoint $SU(5)$ [38, 39, 40, 41] in the case of EPWM.²

Table 1: Quantum numbers of the new particles in EMM under $G_{\text{SM}} \times Z_2$.

Field	N_i	N	η
Z_2	+	−	−
$SU(3)_c$	1	1	1
$SU(2)_L$	1	1	2
$U(1)_Y$	0	0	1/2

2.1 Generation of tree and loop terms

The most general renormalizable CP conserving scalar potential in EMM and EPWM is analogous to the one in the inert doublet model [42, 43, 44]

$$V = \mu_1^2 |\phi|^2 + \mu_2^2 |\eta|^2 + \lambda_1 |\phi|^4 + \lambda_2 |\eta|^4 + \lambda_3 |\phi|^2 |\eta|^2 + \lambda_4 |\phi^\dagger \eta|^2 + \frac{\lambda_5}{2} [(\phi^\dagger \eta)^2 + \text{H.c.}] , \quad (1)$$

²Important for LG is whether the lightest Majorana fermion N_1 is $SU(2)$ singlet or triplet. According to this, in general, both *singlet* and *triplet* types of extensions can be considered for Ma model, and same for Perez–Wise model. However in our definitions EMM (EPWM) generates singlet (triplet) LG.

Table 2: Quantum numbers of the new particles in EPWM and corresponding particles in Adjoint $SU(5)$ under G_{SM} .

Field	$N_1 \equiv \rho_3$	$N_2 \equiv \rho_0$	$N \equiv \rho_8$	$\eta \equiv S_8$
$SU(3)_c$	1	1	8	8
$SU(2)_L$	3	1	1	2
$U(1)_Y$	0	0	0	1/2

where μ_i^2 and λ_i are real, ϕ is the SM Higgs doublet, and the traces over color matrices are assumed in EPWM. The Higgs boson mass squared is $M_h^2 = 4\lambda_1 v_0^2$, where $v_0 = 174$ GeV is the Higgs VEV. The inert doublet η has zero VEV. With positive squared masses of scalars this potential is bounded from below if and only if [45]

$$\lambda_{1,2} > 0, \quad \lambda_3 > -\sqrt{\lambda_1 \lambda_2}, \quad |\lambda_5| < 2\sqrt{\lambda_1 \lambda_2} + \lambda_3 + \lambda_4. \quad (2)$$

The Yukawa and mass terms in the Lagrangian relevant for the neutrino masses can be written as

$$\begin{aligned} -\mathcal{L}_\nu &= Y_{i\alpha} \bar{L}_\alpha N_i \phi + \frac{1}{2} \bar{N}_i M_{ij} N_j^c \\ &+ h_\alpha \bar{L}_\alpha N \eta + \frac{1}{2} \bar{N} M_N N^c + \text{H.c.}, \end{aligned} \quad (3)$$

where α is flavor index, $L = (e_L, \nu_L)^T$ is the SM lepton doublet, and proper contractions of the $SU(2)_L$ and color indexes should be done in EPWM. By integrating out N_i and calculating the loop in Fig. 2 we have

$$\mathcal{L}_\nu^{\text{eff}} = \frac{1}{2} \bar{L} \phi (Y^T M^{-1} Y) \phi^T L + \frac{1}{2\Lambda} \bar{L} \phi (hh^T) \phi^T L + \text{H.c.} \quad (4)$$

In the mass basis of N_i , $M = \text{diag}(M_1, M_2) \equiv D_M$, after absorption of the minus sign by rotation of $\nu_{\alpha L}$ fields the neutrino mass matrix can be written as

$$M_\nu = v_0^2 Y^T D_M^{-1} Y + \frac{v_0^2}{\Lambda} hh^T \equiv M_\nu^{\text{tree}} + M_\nu^{\text{loop}}, \quad (5)$$

where M_ν^{tree} is type I [2, 3, 4, 5, 6, 7] (type I plus III [46, 47]) see-saw contribution in EMM (EPWM), and Λ is the high-energy mass scale, generated in loop, which may be positive or negative, depending on the relevant couplings. For the loop, shown in Fig. 2,

$$\Lambda = \frac{16\pi^2}{C\lambda_5} F^{-1} \left(\frac{M_\eta}{M_N} \right) M_N \simeq \frac{8\pi^2}{C\lambda_5} \left(\ln \frac{M_N}{M_\eta} - \frac{1}{2} \right)^{-1} M_N \quad \text{for } M_\eta \ll M_N, \quad (6)$$

where $C = 1$ in EMM and $C = N_c^2 - 1 = 8$ ($N_c = 3$ is the number of colors) in EPWM, and the loop function

$$\begin{aligned} F(x) &= \frac{x^2 - 1 - \ln x^2}{(1 - x^2)^2} \\ &= -(1 + 2 \ln x) + \mathcal{O}(x^2 \ln x) \quad \text{for } x \ll 1 \end{aligned} \quad (7)$$

comes from the finite part of the Passarino–Veltman function B_0 [48, 49]. We note that the only difference between the loop contributions to the neutrino masses in EMM and EPWM is encoded by the factor C in Eq. (6).

The difference between the new physics contributions of neutral and charged current processes at low energies is measured by the T parameter [1], which is constrained by the present experiments as $T = 0.05 \pm 0.11$ [50]. The contribution of the fermionic triplet to T is zero in the case of mass degeneracy of its neutral and charged components, because the contributions to self-energies of the Goldstones ϕ^+ and χ (see Appendix C in [43] for $\Delta\rho$, which is proportional to T) cancel each other.

2.2 Connection of tree and loop terms

Consider EMM or EPWM as a part of more general theory, which possesses analytical relation among the Yukawa couplings Y and h in Eq. (3). For simplicity, let it be a linear relation

$$h_\alpha^T = a_i Y_{i\alpha} \quad (8)$$

with real a_i . In particular, in Adjoint $SU(5)$ model

$$a_1 = \mp \frac{1}{16\sqrt{6}} \frac{v_0}{|v_{45}|}, \quad a_2 = \pm \frac{\sqrt{5}}{24\sqrt{2}} \frac{v_0}{|v_{45}|}, \quad (9)$$

where v_{45} is the VEV of $\mathbf{45}_H$ representation. Because the same contraction of the representations $\bar{\mathbf{5}}\text{-}\mathbf{24}\text{-}\mathbf{45}_H$ contains the terms, which contribute to both tree and loop level neutrino masses [39]. More explicitly, the term $p_\alpha \bar{\mathbf{5}}_\alpha \mathbf{24} \mathbf{45}_H$ in the Lagrangian, which is involved in the generation of the loop neutrino mass term proportional to $p_\alpha p_\beta$, generates also type I and III see-saw contributions to the neutrino masses, which are dependent on the same coefficient p_α . Notice that in Eq. (9) $|a_i| \lesssim 1$.

Using Eq. (8), Eq. (5) could be rewritten as

$$M_\nu = v_0^2 Y^T \mathcal{M}^{-1} Y, \quad (10)$$

where in the case of two N_i

$$\mathcal{M}^{-1} = \begin{pmatrix} M_1^{-1} + a_1^2 \Lambda^{-1} & a_1 a_2 \Lambda^{-1} \\ a_1 a_2 \Lambda^{-1} & M_2^{-1} + a_2^2 \Lambda^{-1} \end{pmatrix}. \quad (11)$$

The neutrino mass matrix in Eq. (10) can be rewritten as

$$M_\nu = v_0^2 Y'^T D_{\mathcal{M}}^{-1} Y', \quad (12)$$

by using the orthogonal transformation

$$QY = Y', \quad (13)$$

$$Q\mathcal{M}Q^T = D_{\mathcal{M}}, \quad (14)$$

where

$$\mathcal{M} = \frac{M_1 M_2 \Lambda}{\Lambda + a_1^2 M_1 + a_2^2 M_2} \begin{pmatrix} M_2^{-1} + a_2^2 \Lambda^{-1} & -a_1 a_2 \Lambda^{-1} \\ -a_1 a_2 \Lambda^{-1} & M_1^{-1} + a_1^2 \Lambda^{-1} \end{pmatrix} \equiv M_0^2 \begin{pmatrix} a & c \\ c & b \end{pmatrix} \quad (15)$$

is modified mass matrix of heavy fermions, $D_{\mathcal{M}} = \text{diag}(\tilde{M}_1, \tilde{M}_2)$ with the eigenvalues

$$\tilde{M}_{1,2} = \frac{M_0^2}{2} \left(a + b \mp \sqrt{(b-a)^2 + 4c^2} \right) = \frac{M_0^2}{2} \left(a + b \mp \sqrt{(a+b)^2 - 4M_0^{-2}} \right) \quad (16)$$

and

$$Q = \begin{pmatrix} \cos q & \sin q \\ -\sin q & \cos q \end{pmatrix} \quad (17)$$

is real orthogonal matrix with the mixing

$$\sin q = -\frac{\sqrt{2}c}{\sqrt{(b-a)[(b-a) + \sqrt{(b-a)^2 + 4c^2}] + 4c^2}}. \quad (18)$$

Eqs. (15) and (16) show that \tilde{M}_2 has singularity at $\Lambda = -a_1^2 M_1 - a_2^2 M_2$.

For hierarchical N_i with $M_1 \ll \min(|\Lambda|, M_2)$ we have following approximations

$$\tilde{M}_1 \simeq M_1 \left(1 - a_1^2 \frac{M_1}{\Lambda} \right), \quad \tilde{M}_2 \simeq \frac{M_2 \Lambda}{\Lambda + a_2^2 M_2} \quad (19)$$

and

$$\sin q \simeq a_1 a_2 \frac{M_1}{\Lambda}, \quad \cos q \simeq 1 - \frac{a_1 a_2 M_1}{2\Lambda}. \quad (20)$$

2.3 Parametrization of neutrino masses

For explanation of the neutrino experimental data we use the standard Casas-Ibarra [51] parametrization of the Yukawa couplings Y' as

$$Y' = v_0^{-1} D_{\mathcal{M}}^{1/2} \Omega D_{\nu}^{1/2} U^\dagger, \quad (21)$$

where Ω is complex orthogonal (or partly orthogonal for the number of N_i different from three) matrix, and U is the PMNS lepton mixing matrix, which diagonalizes the neutrino mass matrix in the flavor basis according to

$$U^T M_\nu U = \text{diag}(m_1, m_2, m_3) \equiv D_\nu. \quad (22)$$

In the case of two N_i one of the light neutrinos is massless. Hence the quasi-degenerate neutrinos are forbidden, and the only allowed neutrino mass spectra are

- Normal Hierarchical (NH)

$$m_1 = 0, \quad m_2 = \sqrt{\Delta m_{\text{sol}}^2}, \quad m_3 = \sqrt{\Delta m_{\text{atm}}^2}; \quad (23)$$

- Inverted Hierarchical (IH)

$$m_3 = 0, \quad m_1 = \sqrt{\Delta m_{\text{atm}}^2 - \Delta m_{\text{sol}}^2}, \quad m_2 = \sqrt{\Delta m_{\text{atm}}^2}; \quad (24)$$

where $\Delta m_{\text{sol}}^2 = 7.65 \times 10^{-5} \text{ eV}^2$ and $\Delta m_{\text{atm}}^2 = 2.40 \times 10^{-3} \text{ eV}^2$ are the mass-squared differences of solar and atmospheric neutrino oscillations [1]. In this case Ω is 2×3 matrix, which can be written as [52]

$$\Omega^{\text{NH}} = \begin{pmatrix} 0 & \cos z & \pm \sin z \\ 0 & -\sin z & \pm \cos z \end{pmatrix}, \quad \Omega^{\text{IH}} = \begin{pmatrix} \cos z & \pm \sin z & 0 \\ -\sin z & \pm \cos z & 0 \end{pmatrix} \quad (25)$$

in the normal and inverted hierarchy, respectively; where z is the complex angle.

3 Leptogenesis

3.1 CP asymmetry

3.1.1 General formulas

The CP asymmetry is generated in the decays of N_i . Relevant for the unflavored LG total CP asymmetry can be defined as

$$\epsilon_i = \frac{\sum_{\alpha} [\Gamma(N_i \rightarrow e_{\alpha} \phi^{\dagger}) - \Gamma(N_i \rightarrow \bar{e}_{\alpha} \phi)]}{\sum_{\alpha} [\Gamma(N_i \rightarrow e_{\alpha} \phi^{\dagger}) + \Gamma(N_i \rightarrow \bar{e}_{\alpha} \phi)]}. \quad (26)$$

Assuming for the couplings of scalar potential $\max(|\lambda_3|, |\lambda_4|) \ll |\lambda_5|$ to suppress possible two-loop effects, the CP asymmetry can be rewritten as [10, 16, 53, 54]

$$\epsilon_i = \frac{1}{8\pi (YY^{\dagger})_{ii}} \sum_{j \neq 1} \text{Im} \left[(YY^{\dagger})_{ij}^2 \right] f \left(\frac{M_j^2}{M_i^2} \right), \quad (27)$$

where in EMM

$$f(x) = \sqrt{x} \left[\frac{1}{1-x} + 1 - (1+x) \ln \left(\frac{1+x}{x} \right) \right] = -\frac{3}{2\sqrt{x}} + \mathcal{O}(x^{-3/2}) \quad \text{for } x \gg 1, \quad (28)$$

and in EPWM

$$f(x) = \sqrt{x} \left[1 - (1+x) \ln \left(\frac{1+x}{x} \right) \right] = -\frac{1}{2\sqrt{x}} + \mathcal{O}(x^{-3/2}) \quad \text{for } x \gg 1 \quad (29)$$

since the only non-vanishing contribution comes from the vertex correction [40].

The decay parameter can be written as

$$K \equiv \frac{\tilde{\Gamma}_D}{H|_{T=M_{\rho_3}}} = \frac{\tilde{m}}{m_*}, \quad (30)$$

where $\tilde{\Gamma}_D$ is equal to the total decay rate Γ_D in EMM and $\tilde{\Gamma}_D = \Gamma_D/3$ in EPWM, where it is normalized by the number of components of the triplet Majorana fermion. The rescaled decay rate (effective neutrino mass) is defined as [55]

$$\tilde{m} \equiv 8\pi \frac{v_0^2}{M_1^2} \tilde{\Gamma}_D = \frac{v_0^2}{M_1} (YY^\dagger)_{11}, \quad (31)$$

and the rescaled Hubble expansion rate (equilibrium N_1 mass) is

$$m_* \equiv 8\pi \frac{v_0^2}{M_1^2} H|_{T=M_1} \simeq 1.08 \times 10^{-3} \text{ eV}. \quad (32)$$

For NH (IH) the strong washout regime requires

$$K \geq K_{\text{sol(atm)}} \equiv m_{2(1)}/m_* \simeq 8.1 \text{ (46)} \gg 1. \quad (33)$$

3.1.2 Hierarchical N_i

In the hierarchical limit $M_1/M_{i>1} \rightarrow 0$, Eq. (27) can be rewritten as

$$\epsilon_1 = -\frac{A}{(YY^\dagger)_{11}} \sum_{j \neq 1} \frac{M_1}{M_j} \text{Im} \left[(YY^\dagger)_{1j}^2 \right] = -\frac{AM_1}{(YY^\dagger)_{11}} \Sigma \quad (34)$$

with $A = 3/(16\pi)$ and $1/(16\pi)$ in the EMM and EPWM, respectively, and

$$\Sigma \equiv \sum_{j \neq 1} \text{Im} \left[(YY^\dagger)_{1j}^2 M_j^{-1} \right] = \sum_{j=1,2,\dots} \text{Im} \left[(YY^\dagger)_{1j}^2 M_j^{-1} \right]. \quad (35)$$

Using Eqs. (13) and (21), we have

$$(YY^\dagger)_{11} = \frac{1}{v_0^2} \left(Q^T D_{\mathcal{M}}^{1/2} \Omega D_\nu \Omega^\dagger D_{\mathcal{M}}^{1/2} Q \right)_{11} = \frac{1}{v_0^2} \sum_\alpha m_\alpha \left| Q^T D_{\mathcal{M}}^{1/2} \Omega \right|_{1\alpha}^2. \quad (36)$$

Using Eq. (5), we get

$$\Sigma = \text{Im} \left[(YY^\dagger D_M^{-1} Y^* Y^T)_{11} \right] = \frac{1}{v_0^2} \text{Im} \left\{ [Y(M_\nu - M_\nu^{\text{loop}})^\dagger Y^T]_{11} \right\} \equiv \Sigma_\nu + \Sigma_\nu^{\text{loop}}, \quad (37)$$

where Σ_ν (Σ_ν^{loop}) is the term with M_ν (M_ν^{loop}). Using Eqs. (13), (21) and (22), Σ_ν can be rewritten as

$$\Sigma_\nu = \frac{1}{v_0^4} \text{Im} \left[\left(Q^T D_{\mathcal{M}}^{1/2} \Omega D_\nu^2 \Omega^T D_{\mathcal{M}}^{1/2} Q \right)_{11} \right] = \frac{1}{v_0^4} \sum_\alpha m_\alpha^2 \text{Im} \left[\left(Q^T D_{\mathcal{M}}^{1/2} \Omega \right)_{1\alpha}^2 \right], \quad (38)$$

and, using Eqs. (5) and (8), Σ_ν^{loop} can be rewritten as

$$\Sigma_\nu^{\text{loop}} = -\frac{1}{\Lambda} \text{Im} [(Yh^*)_1^2] = -\frac{1}{\Lambda} \text{Im} [(YY^\dagger a^\dagger)_1^2]. \quad (39)$$

In the case of two N_i we have

$$\Sigma = \frac{1}{M_2} \text{Im} [(YY^\dagger)_{12}^2] \quad (40)$$

and

$$\Sigma_\nu^{\text{loop}} = -\frac{a_2^2}{\Lambda} \text{Im} [(YY^\dagger)_{12}^2] = -a_2^2 \frac{M_2}{\Lambda} \Sigma. \quad (41)$$

From Eqs. (37), (38) and (41), we get

$$\Sigma = \frac{M'_2}{M_2} \Sigma_\nu = \frac{1}{v_0^4} \frac{M'_2}{M_2} \sum_\alpha m_\alpha^2 \text{Im} \left[\left(Q^T D_{\mathcal{M}}^{1/2} \Omega \right)_{1\alpha}^2 \right], \quad (42)$$

with

$$M'_2 = \left(\frac{1}{M_2} + \frac{a_2^2}{\Lambda} \right)^{-1}, \quad (43)$$

and, using Eqs. (34) and (36), we have

$$\epsilon_1 \simeq -A\mu \frac{M_1}{v_0^2} \frac{\sum_\alpha m_\alpha^2 \text{Im} \left[\left(Q^T D_{\mathcal{M}}^{1/2} \Omega \right)_{1\alpha}^2 \right]}{\sum_\alpha m_\alpha \left| Q^T D_{\mathcal{M}}^{1/2} \Omega \right|_{1\alpha}^2}. \quad (44)$$

We note that the magnification factor $\mu \equiv M'_2/M_2$ is formally equivalent to the magnification of thin concave lens with the focal length $f = -\Lambda/a_2^2$ since Eq. (43) can be rewritten as

$$\frac{1}{M_2} - \frac{1}{M'_2} = \frac{1}{f}. \quad (45)$$

3.2 Boltzmann equations

Boltzmann equations in the unflavoured regime can be written as (for more details see [39, 40, 41, 55, 56] and Refs. therein)

$$\frac{dN_{N_1}}{dz} = -(D + S)(N_{N_1} - N_{N_1}^{\text{eq}}), \quad (46)$$

$$\frac{dN_{B-L}}{dz} = -\epsilon_1 D(N_{N_1} - N_{N_1}^{\text{eq}}) - W N_{B-L}, \quad (47)$$

where $z = M_1/T$, and N_X ($X = N_1, B - L$) is the number density of X calculated in a co-moving volume containing one N_1 (all of its components) in ultrarelativistic thermal equilibrium: $N_{N_1}^{\text{eq}}(T \gg M_1) = 1$. Initially, $N_{B-L}^{\text{eq}}(T \gg M_1) = 0$. $D = \Gamma_D/(Hz)$ is the decay factor, and W is the washout term. The scattering term is $S = S_\phi$ in EMM and $S = S_\phi + 2S_g(N_{N_1} + N_{N_1}^{\text{eq}})$ in EPWM (see footnote on p. 3), where S_ϕ is the contribution from Higgs-mediated scatterings, and the gauge scattering of the triplet Majorana fermion [16] can be fitted by [40]

$$S_g \simeq 10^{-3} \frac{M_{\text{Pl}}}{M_{\rho_3}} \frac{\sqrt{1 + \pi z^{-0.3}/2}}{(15/8 + z)^2(1 + \pi z/2)} e^{0.3z}, \quad (48)$$

where $M_{\text{Pl}} = 1.221 \times 10^{19}$ GeV is the Planck mass.

After solving the system of Boltzmann equations (46)–(47), we obtain $N_{B-L}^{\text{f}} = N_{B-L}(z \rightarrow \infty)$ (in the calculations below the final value $z = 10$ is used, where the fit in Eq. (48) is still applicable), included in the final baryon asymmetry

$$\eta_B \simeq 3 \times 0.88 \times 10^{-2} N_{B-L}^{\text{f}}. \quad (49)$$

This result should be compared with the allowed values

$$5.1 \times 10^{-10} < \eta_B^{\text{BBN}} < 6.5 \times 10^{-10}, \quad (50)$$

which come from the nucleosynthesis predictions and observed abundances of light elements [1].

3.3 Analysis

3.3.1 Non-resonant case

For the particular spectrum $M_1 \ll \min(|\Lambda|, M_2, |\Lambda + a_2^2 M_2|) \equiv M_{\text{min}}$, using Eqs. (19) and (20), Eqs. (36) and (38) can be rewritten as

$$(YY^\dagger)_{11} = \frac{\tilde{M}_1}{v_0^2} \left[Q_{11}^2 \sum_\alpha m_\alpha |\Omega_{1\alpha}|^2 + \mathcal{O} \left(\sqrt{\frac{M_1}{M_{\text{min}}}} \right) \right], \quad (51)$$

$$\Sigma_\nu = \frac{\tilde{M}_1}{v_0^4} \left\{ Q_{11}^2 \sum_\alpha m_\alpha^2 \text{Im} [(\Omega_{1\alpha})^2] + \mathcal{O} \left(\sqrt{\frac{M_1}{M_{\text{min}}}} \right) \right\}. \quad (52)$$

Hence the CP asymmetry in Eq. (44) can be rewritten as

$$\epsilon_1 \simeq -A\mu \frac{M_1}{v_0^2} \frac{\sum_\alpha m_\alpha^2 \text{Im} (\Omega_{1\alpha}^2)}{\sum_\alpha m_\alpha |\Omega_{1\alpha}|^2} = -A\mu \frac{M_1}{v_0^2} \frac{(m_b^2 - m_a^2) \text{Im} \sin^2 z}{m_a |\cos z|^2 + m_b |\sin z|^2}, \quad (53)$$

where $a = 2$ (1) and $b = 3$ (2) for NH (IH) neutrinos. Eq. (53) results in the upper bound for the CP asymmetry

$$|\epsilon_1| \lesssim A\mu \frac{M_1}{v_0^2} (m_b - m_a), \quad (54)$$

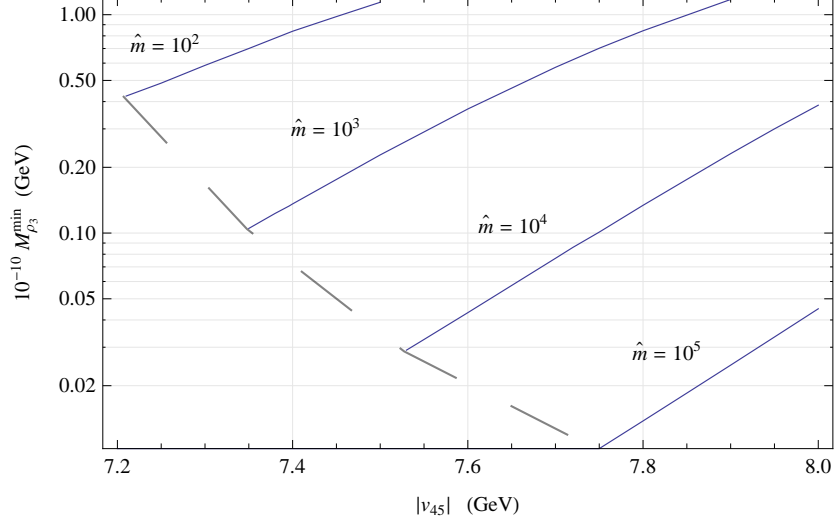


Figure 3: Minimal allowed by unflavoured LG values of M_{ρ_3} vs. $|v_{45}|$ for $K = 20$, $\lambda_5 = -1/2$ and NH neutrinos in Adjoint $SU(5)$. (Chosen values of \hat{m} cover the region allowed by unification at 1-loop level, proton decay and collider searches [57].)

which is equal to the DI bound [14] (its non-supersymmetric version has the same factor A as EMM), rescaled by μ . Using Eq. (51), Eq. (31) can be rewritten as

$$\tilde{m} \simeq Q_{11}^2 \frac{\tilde{M}_1}{M_1} \sum_{\alpha} m_{\alpha} |\Omega_{1\alpha}|^2 \simeq m_a |\cos z|^2 + m_b |\sin z|^2 \geq m_a, \quad (55)$$

which is the usual form. In the considered non-resonant region the DI bound can not be significantly relaxed since $\mu \lesssim 1$. However the new allowed parameter ranges for successful LG appear for large values of the decay parameter K and for IH neutrino masses, as was shown for triplet LG in [39] (in the context of Adjoint $SU(5)$), using precise formulas for \tilde{m} and ϵ_1 in the case of two N_i .

3.3.2 Resonant case

For the case of approximate cancellation of the tree and loop contributions to the neutrino masses, namely $\Lambda \simeq -a_2^2 M_2$ (see Eqs. (5) and (8)), the factor μ is large and enhances the CP asymmetry in Eqs. (44) and (53).

In Adjoint $SU(5)$ model for $M_{S_8} \equiv M_{\eta} = 1$ TeV and $\lambda_5 = -1/2$ this resonance happens at $|v_{45}| \sim 8$ GeV. By choosing values of $|v_{45}|$ near 8 GeV and using the method of calculations described in [39], we get minimal allowed by successful LG values of $\rho_3 \equiv N_1$ mass versus $|v_{45}|$, shown in Fig. 3 for $K = 20$, NH neutrinos, SM Higgs mass

$M_h = 130$ GeV and several chosen values of $\hat{m} \equiv M_{\rho_8}/M_{\rho_3} \equiv M_N/M_{N_1}$ ³. Below the dashed line in Fig. 3 the unflavoured LG is not allowed. Clearly, for stronger hierarchy of M_i (higher values of \hat{m}) the lower bound on M_{ρ_3} is weaker. Fig. 3 shows that the allowed values of M_{ρ_3} can be lowered by several orders of magnitude comparing to the scale of 10^{11} GeV, which is relevant for the case of vanishing loop contribution to the neutrino masses, see [39, 40].

In the case of singlet LG (as in EMM) the lower bound on the strongly hierarchical heavy neutrino masses (e.g., $M_2/M_1 \gtrsim 10^7$) can be decreased up to the TeV scale, which is testable at the LHC [30, 31, 32, 33, 34, 35, 36]. However it requires fine tuning of the parameters of the theory. We remark that this bound holds for the ordinary right-handed neutrinos in contrast to the reduced by the loop factor $16\pi^2$ DI bound on the masses of Z_2 odd Majorana fermions (N) derived in [58].

We note that the inert doublet model, which is embedded in the considered theory, may provide contribution to the dark matter in the universe, see [44] for recent study.

4 Summary

The SM extensions, which change the usual connection of the leptogenesis to the observable neutrino masses and relax the Davidson-Ibarra bound, are introduced. The lower bound on the hierarchical masses of heavy Majorana fermions can be significantly decreased, while the upper bound on the light neutrino masses may be increased in this theory, which may be tested in the near future experiments. The non-SM particles, involved in the loop contribution to the neutrino masses in Fig. 2, such as scalar octet in Adjoint $SU(5)$ model can be tested at the Large Hadron Collider and next colliders [59, 60, 61]. Finally, the long standing gravitino problem can be solved in the supersymmetric version of this theory.

Acknowledgments

The author thanks Riccardo Barbieri, Kristjan Kannike and Anatoly Borisov for useful discussions and comments, and the organizers of the BLV2011 Workshop Pavel Fileviez Perez and Yuri Kamyshev for hospitality in Gatlinburg. This work was supported in part by the EU ITN “Unification in the LHC Era”, contract PITN-GA-2009-237920 (UNILHC) and by MIUR under contract 2006022501.

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³We note that aside from the resonance area the maximal allowed value of \hat{m} for unflavoured LG is $\sim 10^4$ [39]. However this value can be increased in the resonant case considered here.

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